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Electroweak precision data, light sleptons and stability of the SUSY scalar potential

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Abstract

The light slepton–sneutrino scenario with nonuniversal scalar masses at the GUT scale is preferred by the electroweak precision data. Though a universal soft breaking mass at or below the Plank scale can produce the required nonuniversality at the GUT scale through running, such models are in conflict with the stability of the electroweak symmetry breaking vacuum. If the supergravity motivated idea of a common scalar mass at some high scale along with light sleptons is supported by future experiments that may indicate that we are living in a false vacuum. In contrast $SO(10)$ D-terms, which may arise if this GUT group breaks down directly to the Standard Model, can lead to this spectrum with many striking phenomenological predictions, without jeopardizing vacuum stability.

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The electroweak precision (EWP) tests by the experiments at LEP and SLC [1] are on the whole in excellent agreement with the Glashow–Weinberg–Salam standard model (SM). However, if some judiciously chosen subset of the data is examined, a few unsatisfactory features of the SM fit are revealed [1,2]

- The measured values of the parameter $\sin^2\theta_{\text{eff}}$ from the observables A_{LR} and A^b_{FB} differ at 3.5σ level.
- Moreover, the value of this parameter as given by the hadronic asymmetries and the leptonic asymmetries also exhibit a considerable discrepancy (at the 3.6σ level).
- When a global fit is performed a $\chi^2/\text{d.o.f.} = 26/15$ corresponding to a C.L. $= 0.04$ is obtained, which is hardly satisfactory.
- If the hadronic data is excluded from the global fit the quality of the fit improves considerably ($\chi^2/\text{d.o.f.} = 2.5/3$, corresponding C.L. $= 0.48$) while the exclusion of the leptonic data worsens the fit to an unacceptable level ($\chi^2/\text{d.o.f.} = 15.3/3$, corresponding C.L. $= 0.0016$).

These observations tempt one to conclude that the hadronic data may be plagued by some hitherto

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unidentified experimental problem and, hence, the leptonic data should be taken more seriously [2].

This conclusion is challenged by the direct lower bound on the Higgs mass $m_H > 113$ GeV [3] and its indirect determination from EWP data considering the leptonic asymmetries only [2,4]. Using $\sin^2\theta_{\text{eff}}$ measured from both hadronic and leptonic asymmetries, the central value of the fitted Higgs mass and the 95% C.L. upper limit on it happens to be 98 GeV and 212 GeV, respectively [1]. These values consistent with the direct search limit, have been confirmed by [2]. However, if $\sin^2\theta_{\text{eff}}$ from leptonic data only is employed, the corresponding numbers become 42 GeV and 109 GeV, a situation which is hardly acceptable vis-a-vis the direct limit.

It must be admitted that there are uncertainties in the fitted value of m_H [2]. The result has some sensitivity on the value of $\alpha_{\text{QED}}(m_Z)$ which is scheme dependent although most of the existing schemes lead to upper bounds on m_H in conflict with the direct search limit. Uncalculated higher order effects may have a modest impact on the fitted value of m_H [2]. Finally, if the current 1σ upper limit of the top mass ($m_t = 174.3 + 5.1$ GeV) rather than its central value is used in the fit, then the compatibility of the fitted value of m_H with the direct search limit improves.

Although these uncertainties may conspire to produce an agreement between the leptonic EWP data and the direct limit on m_H within the framework of the SM, the situation is sufficiently provoking to reanalyse the data in extensions of the SM.

One interesting possibility is to extend the discussion within the framework of supersymmetry [5]. Altarelli et al. [2] have found the MSSM parameter space (PS) where the SUSY corrections to the electroweak observables are sufficiently large and act in the direction of improving the quality of fit. The most significant loop contributions come from the sneutrino ($\tilde{\nu}$), in particular, if sneutrino mass is in the in the range 55–80 GeV, and a perfect agreement with the data is obtained with $m_H = 113$ GeV. The charged left-slepton ($\tilde{l}_L$) mass is related to $m_{\tilde{\nu}}$ by the SU(2)-breaking D-term: $m^2_{\tilde{\ell}_L} = m_{\tilde{\nu}_\ell}^2 - \frac{1}{2}m_Z^2\cos 2\beta$, in a model-independent way. Since it must be heavier than 96 GeV according to the LEP direct search limits on charged sleptons [6], the parameter $\tan\beta$ must be moderately large which is not a severe restriction.

This spectrum, however, is incompatible with the popular mSUGRA [7] scenario with a common scalar mass m_0 at the GUT scale (M_G). Within the framework of mSUGRA such light sneutrinos automatically demand even lighter right-sleptons, which are already ruled out by the LEP mass limits on charged sleptons. Thus one has to look for alternatives with nonuniversal scalar masses at M_G. In this Letter we shall look for such alternatives and scrutinize them in the light of vacuum stability.

We shall consider only those class of models where the sfermions of the first two generations are nearly degenerate with mass m_0 at M_G, as is required by the absence of flavour changing neutral currents. Moreover, we shall assume a universal gaugino mass $m_{1/2}$ at M_G as this assumption is likely to be valid in all GUT models irrespective of the specific choice of the GUT group. Given these parameters the left-slepton and sneutrino masses of the first two generation at the weak scale can be computed by using the standard one loop renormalisation group (RG) equations. Other SUSY parameters may influence the running at the two loop level. Using ISAJET-ISASUSY we have convinced ourselves that these higher order corrections are indeed negligible. We constrain m_0 and $m_{1/2}$ by requiring 55 GeV $< m_{\tilde{\nu}} <$ 80 GeV at the weak scale (Fig. 1). The only other relevant SUSY parameter that enters the analysis through the $SU(2)$ breaking D-term is $\tan\beta$, although the dependence on it is rather weak. Almost identical allowed PS is obtained for all $\tan\beta \gtrsim 5$. As long as $\tan\beta$ is not too large (say, $\tan\beta \lesssim 20$), $\tilde{\tau}_L$ will be degenerate with the sleptons of the first two generations (to a very good approximation). For larger $\tan\beta$, it may be somewhat lighter. Since the experimental bound on the $\tilde{\tau}_L$ mass is considerably weaker ($m_{\tilde{\tau}} > 68$ GeV) than that for the selectron and smuon, higher values of $\tan\beta$ can also be considered in principle, although we shall not pursue this case further.

The range of m_0 and $m_{1/2}$ shown in Fig. 1 may be moderately altered if one considers a large hierarchy among the scalar masses at M_G. This happens due to the presence of a particular term in the RG equation which is usually neglected in the mSUGRA approximation (see Eq. (4) and the discussions following it). We shall consider below a specific model with this feature.

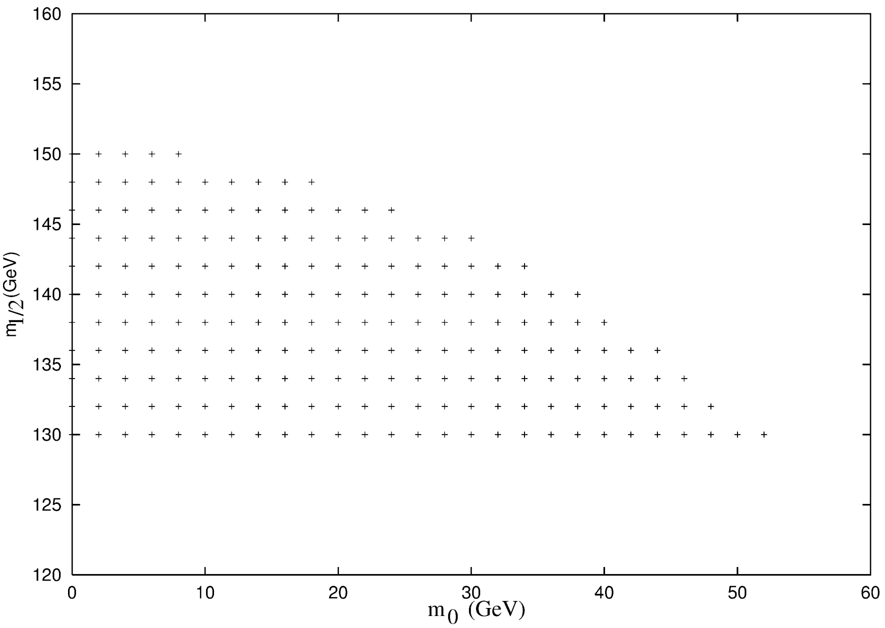

Fig. 1. The APS in the m_0–$m_{1/2}$ plane for 55 GeV $< m_{\tilde{\nu}} <$ 80 GeV with $\tan\beta = 15$. The lower limit on $m_{1/2}$ is due to the chargino mass bound from LEP.

So far no assumption about the other soft breaking parameters was necessary. However, in order to take into account the chargino mass bound $m_{\tilde{\chi}^\pm} >$ 100 GeV [6] and to test the stability of the scalar potential [8,9] one has to specify more SUSY parameters. In general, $m_{\tilde{\chi}^\pm}$ depends on the Higgsino mass parameter (μ) and $\tan\beta$ in addition to $m_{1/2}$. The entire range of $m_{1/2}$ in Fig. 1 is such that μ can be chosen so as to make $m_{\tilde{\chi}^\pm}$ consistent with the LEP bound. Of course, $m_{\tilde{\chi}^\pm}$ is not a very sensitive function of μ unless it is very small ($\mu \lesssim 100$ GeV). We next turn our attention on $m_{\tilde{e}_R}$ and the stability of the scalar potential [8,9].

Before looking into specific models it is worthwhile to focus on some generic features of models with light sleptons. In several recent works [9–11] on the stability of the standard electroweak symmetry breaking (EWSB) vacuum, it has been found that low mass sleptons (to be more specific, sleptons significantly lighter than the electroweak gauginos) are somewhat disfavoured. In view of the fact that there is already a strong lower bound on the chargino mass it is important to check the compatibility of the light sneutrino scenario favoured by the EWP data and vacuum stability.

The unbounded from below 3 (UFB3) direction [9] of the scalar potential, its evaluation procedure and the choice of the generation indices (i, j) which leads to the strongest constraint are elaborately discussed in [9,10]. To clarify why light sleptons are strongly disfavoured, Eq. (93) of [9] has to be examined. The required equation is

$$V_{\mathrm{UFB3}} = \left[m_{H_u}^2 + \tilde{m}_{\tilde{\ell}_{Li}}^2\right]|H_u|^2 + \frac{|\mu|}{\lambda_{e_j}}\left[m_{\tilde{\ell}_{Lj}}^2 + m_{\tilde{\ell}_{Rj}}^2 + m_{\tilde{\ell}_{Li}}^2\right]|H_u| - \frac{2m_{\tilde{\ell}_{Li}}^4}{g_1^2 + g_2^2} \tag{1}$$

with $i \neq j$. Here λ_{e_j} is a leptonic Yukawa coupling and g_1 and g_2 are the $U(1)_Y$ and $SU(2)$ gauge couplings, respectively. The UFB3 constraint arises from the requirement that V_{UFB3} must be shallower than the EWSB minima ($V_{0\,\mathrm{min}}$) (see Eq. (92) of [9]). To get the strongest constraints $i = 1$ and $j = 3$ is considered. Over a large region of the PS corresponding to light sleptons, the first term of Eq. (1) dominates when λ_τ is substituted in the second term. The parameters are evaluated at a judiciously chosen renormalisation scale $\widehat{Q}$, where higher-order loop corrections to the scalar potential are small and may be neglected [9, 12]. At this scale, the mass parameter $m^2_{H_u}$ (H_u refers to the Higgs bosons coupling to the up-type quarks) gets a large negative value which is required by radiative electroweak symmetry breaking (REWSB). Thus the first term tends to violate the UFB3 constraint for small values of $m^2_{\tilde{\ell}_{Li}}$. In fact, it has been shown in reference [11] that the anomaly mediated supersymmetry breaking (AMSB) model with light sleptons violate the UFB3 constraint.

We now wish to scrutinize the PS favoured by EWP data (Fig. 1) in the light of the stability of the vacuum. At this stage we have to be more specific about the model since $m^2_{H_u}$, $m^2_{\tilde{L}_{Rj}}$ and $|\mu|$ are also needed to check this point. We first consider a SU(5) SUSY GUT with a common scalar mass m_0 at the Plank scale ($M_P \approx 2 \times 10^{18}$ GeV) [13] instead of M_G. An attractive feature of this model is that for the first two generations the mass of $\tilde{l}_R$ (denoted by m_{10} at M_G) belonging to the 10-plet of SU(5) happens to be larger than that of left slepton belonging to the $\bar{5}$ representation (denoted by m_5 at M_G) due to the running between M_P and M_G. Thus the conflict between the low mass sneutrino and the LEP limit on $\tilde{\ell}_R$ mass seems to be resolved, at least qualitatively.

For the 3rd generation, m_{10} may be somewhat smaller if the relevant Yukawa couplings happen to be large at M_G and contribute to the running (all relevant RG equations are given in Ref. [13]). This, however, may not be a serious problem since the limit on $m_{\tilde{\tau}_R}$ is considerably weaker as discussed above.

When we look into the numerical details the situation, however, is far from simple. According to Polansky et al. the GUT scale values m_{10} and m_5 for the first two generations are approximately [13]

$$m^2_{10} = m^2_0 + 0.45\, m^2_{1/2}, \tag{2}$$

$$m^2_5 = m^2_0 + 0.30\, m^2_{1/2}, \tag{3}$$

assuming that SUGRA generates the common scalar mass m_0 exactly at M_P. Since $m_{1/2}$ has to be greater than 130 GeV (approximately; see Fig. 1) we find that m_5 is too large to give $m_{\tilde{\nu}}$ in the required range at the weak scale even if $m_0 \approx 0$. We note that if the common soft breaking mass is generated well below the Plank scale this difficulty may be avoided. Moreover, GUT threshold corrections, which cannot be computed precisely without specifying other GUT parameters like masses of heavy multiplets, may affect both m_{10} and m_5 to some extent. In view of these uncertainties one cannot discard this model on this ground alone. We shall henceforth treat m_{10} and m_5 as phenomenological parameters at M_G with $m_{10} > m_5$. Their actual values are to be chosen such that all charged slepton masses at the weak scale satisfy the LEP bound.

The Achilles' heel of the model however, happens to be the running of $m^2_{H_u}$ between M_P and M_G. This running is controlled by not only the Yukawa couplings h_t and h_b but also by λ the coupling of the scalars belonging to the 5-, $\bar{5}$- and 24-plet of $SU(5)$. In course of running $m^2_{H_u}$ is usually reduced as one goes below M_P, whereas m_5 driven by the gauge coupling alone increases. After considering various scenarios with different magnitudes of these couplings Ref. [13] has concluded that $m^2_{H_u} \lesssim m_5$ in general, while the equality holds if all the Yukawa couplings and λ are negligibly small. We have checked that in such scenarios the UFB3 constraint is always violated for the PS in Fig. 1 as is suggested by Eq. (1).

Of course, moderate shifts in $m^2_{H_u}$ and m_5 may come from GUT threshold corrections [13] which may lead to $m^2_{H_u} > m_{\bar{5}}$. The magnitude of this shift depends on the details of the GUT model and we do not attempt to compute it. However, adjusting m_5 and m_{10} such that both $\tilde{l}_L$ and $\tilde{l}_R$ satisfy the experimental bounds at the weak scale, we find that $m_{H_u}, m_{10} \gg m_5$ is needed to satisfy the UFB3 constraint (see Table 1 for sample values). Such large splittings between m_5 and other GUT scale masses is unlikely to arise from threshold corrections.

If one considers an $SO(10)$ SUSY GUT instead, the matter fields of the first two generations belonging to the 16-plet remain degenerate at M_G even if running below M_P is considered. This will inevitably lead to

Table 1
Sample GUT scale masses and consistency with the UFB3 condition for $A_0 = 0$, $\tan\beta = 15$, $\mu > 0$, $m_{1/2} = 153$ GeV

$m_{\tilde{e}_L}$ (GeV)	$m_{\tilde{e}_R}$ (GeV)	m_{H_u} (GeV)	m_{H_d} (GeV)	
36	210	240	340	UFB3 allowed
36	210	220	340	UFB3 disallowed
36	210	240	300	UFB3 disallowed

a light $\tilde{l}_R$ at the weak scale if the sneutrino mass is required to be in the range preferred by EWP data.

Thus running above the GUT scale alone in a SUGRA type scenario with a common scalar mass generated between M_P and M_G, is not likely to yield the spectrum preferred by EWP data if the stability of the vacuum is taken into account.

If one gives up the UFB3 constraint by assuming that the standard vacuum is only a false vacuum [14], while the global minimum of the scalar potential is indeed charge color breaking then the above constraints do not apply. If the tunnelling time for transition between the false vacuum and the true vacuum happens to be much larger than the age of the universe, such a model cannot be rejected outright, although it seems to be against our intuitive notion of stability. Moreover, the tunnelling time, which can be routinely calculated in models with a single scalar, cannot be computed reliably in models with multiple scalars. Yet the conclusions derived in the preceding paragraphs do not loose their significance. If future experimental data confirms light sleptons along with a mass spectrum stemming from a SUGRA motivated common scalar mass at some high scale $\lesssim M_P$, then that would indicate that we may be living in a false vacuum, no matter how counter intuitive it may appear to be at the first sight.

The remaining of this paper shall deal with a type of nonuniversality which arises when a GUT group breaks down to a group of lower rank leading to nonuniversal D-terms at M_G [15]. This type of models can produce the spectrum preferred by EWP data without violating the UFB3 constraint. As a specific example we consider an $SO(10)$ SUSY GUT breaking down to the SM in a single step. The relevant mass formulae at M_G are:

$$m^2_{\tilde{Q}} = m^2_{\tilde{E}} = m^2_{\tilde{U}} = m^2_{16} + m^2_D,$$
$$m^2_{\tilde{D}} = m^2_{\tilde{L}} = m^2_{16} - 3m^2_D,$$
$$m^2_{H_{d,u}} = m^2_{10} \pm 2m^2_D,$$

where m_D is the D-term with unknown magnitude, the common mass of all the members of the 16-plet of $SO(10)$ at M_G is denoted by m_{16} and the common Higgs mass by m_{10}.

This model is interesting since even though all sfermion masses are degenerate at M_G, which indeed should be the case for the first two generations of sfermions as discussed above, the D-terms may introduce significant nonuniversality between the L and R sleptons making the latter somewhat heavier than the former. Thus a light sneutrino as required by the EWP data does not necessarily imply a lighter R-slepton.

In general, the Higgs mass m_{10} and m_{16} could be different at M_G due to the running between M_P and M_G. However, it is interesting to note that even if m_{10} and m_{16} are nearly degenerate at M_G, the D-term may make $m^2_{H_u}$ significantly heavier that the left sleptons at M_G. Because of this reason the model can be UFB3 stable without requiring m_{10} to be much larger than m_{16}. We shall consider both universal ($m_{16} = m_{10}$) and nonuniversal ($m_{16} \neq m_{10}$) scenarios.

The methodology of finding the spectra is same as in [10]. μ and B are determined by the REWSB condition at a scale $M_S = \sqrt{m_{\tilde{t}_L} m_{\tilde{t}_R}}$. Then we put the experimental constraints. For a given m_{16} and $m_{1/2}$, we consider the smallest m_D such that $m_{\tilde{\nu}} <$ 80 GeV. Larger values of m_D may also be considered provided $m_{\tilde{\nu}}$ is in the range 55 GeV $< m_{\tilde{\nu}} <$ 80 GeV. However, larger values of m_D tends to yield stronger UFB3 constraints.

We first discuss the APS without requiring Yukawa unification, in the m_{16}–$m_{1/2}$ plane for $m_{16} = m_{10}$, $A_0 = 0$, $\tan\beta = 15$ and $\mu > 0$ as shown in Fig. 2. The upper bound on $m_{1/2}$ for a given m_{16} corresponds to the situation when no m_D can give $m_{\tilde{\nu}_{e,\mu}} \leqslant 80$ GeV and the lower bound by experimental lower limit on $\tilde{\chi}^{\pm}$. The D-term can control $m_{\tilde{\ell}_L}$ and, hence $m_{\tilde{\nu}}$, over a large range of m_{16}, which, therefore, is found to be large. If we increase m_{16} further, the contribution from τ Yukawa coupling decreases $m_{\tilde{\tau}_L}$ even for $\tan\beta = 15$ thanks to a large $m_{\tilde{\ell}_R}$. As a result $m_{\tilde{\nu}\tau}$ falls below the experimental bound (43.6 GeV), even though $m_{\tilde{\nu}_{e,\mu}}$ are in the vicinity of 80 GeV. The upper and lower limits on m_{16} significantly depends on A_0 and $\tan\beta$.

The fact that the allowed range of $m_{1/2}$ increases with m_{16} is rather puzzling. The origin of this lies in a term in the RG equation which is usually neglected in mSUGRA.

$$\frac{dm_{\tilde{\ell}_L}}{dQ} = \frac{3}{8\pi^2}\big[-0.6g_1^2M_1^2 - 3g_2^2M_2^2$$
$$- 0.3g_1^2\{m^2_{H_u} - m^2_{H_d} + \big(2m^2_{\tilde{u}_L} + m^2_{\tilde{t}_L}\big)$$

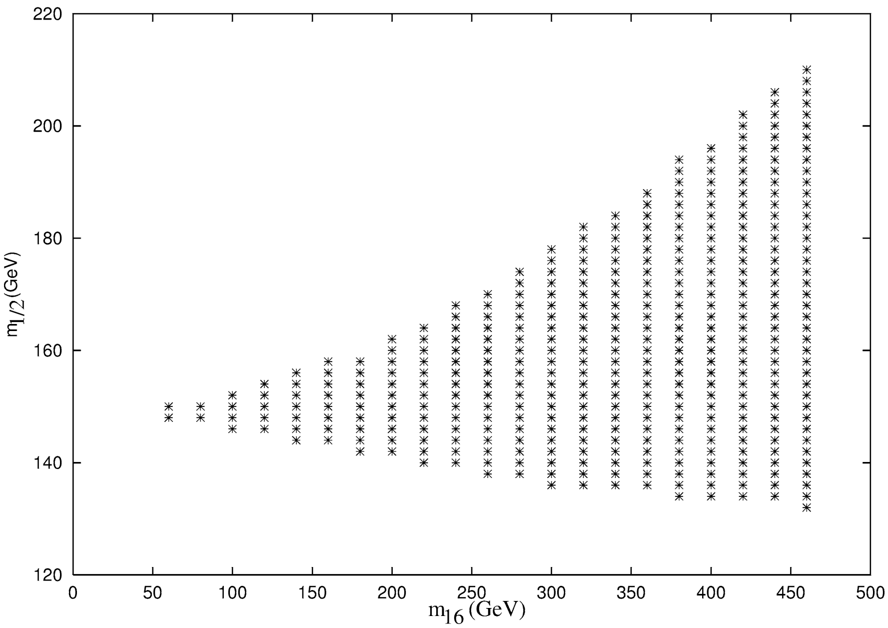

Fig. 2. The APS for 55 GeV $< m_{\tilde{\nu}} <$ 80 GeV in the $SO(10)$ model with $m_{10} = m_{16}$, $A_0 = 0$, $\tan\beta = 15$ and sign$(\mu) > 0$ and m_D is fixed by the light sleptons criterion. In our notation a $*$ denotes a point ruled out by UFB3 while a $+$ indicates a UFB3 allowed point.

$$- \left(2m^2_{\tilde{e}_L} + m^2_{\tilde{\tau}_L}\right) - 2\left(2m^2_{\tilde{u}_R} + m^2_{\tilde{t}_R}\right) + \left(2m_{\tilde{d}_R} + m_{\tilde{b}_R}\right) + \left(2m_{\tilde{e}_R} + m^2_{\tilde{\tau}_R}\right)\big\}\big]. \quad (4)$$

The last term on the right-hand side is zero at M_G in the mSUGRA model. Moreover, its coefficient is rather small. Hence, the contribution of this term remains small even at the weak scale. In the D-term model, however, this term is already large at the GUT scale in particular due to the $m^2_{H_u} - m^2_{H_d}$ term. This difference is indeed large if the D-term is chosen to be large in order to have $m_{\tilde{\nu}}$ in the desired range. The slepton and sneutrino masses are reduced under the influence of this term by as much at 10–15 GeV for large m_{16}. As a result unexpectedly large values of $m_{1/2}$ can be accommodated.

If $\tan\beta$ is lowered, the mass of lightest Higgs (m_h) decreases rapidly, low values of m_{16} are not allowed if $m_h \gtrsim 113$ GeV is required. However, if m_{16} is increased, the Higgs mass increases appreciably through radiative corrections. Moreover, the running of $m_{\tilde{\tau}_L}$ and hence of $m_{\tilde{\tau}_\nu}$, are also modest for low $\tan\beta$. Due to these reasons higher values of m_{16} are allowed. We find 300(60) GeV $\lesssim m_{16} \lesssim$ 700(460) GeV for $\tan\beta = 7(15)$, while the other parameters are the same as in Fig. 2.

Increasing the absolute value of A_0 makes large difference between $m_{\tilde{\nu}e,\mu}$ and $m_{\tilde{\nu}\tau}$. As a result m_{16} gets a stringent upper bound. It also lowers m_H very rapidly giving a strong lower bound on m_{16}. For example, $60(120) \lesssim m_{16} \lesssim 460(420)$ GeV for $A_0 = 0(m_{16})$, the other parameters being the same as in Fig. 2.

There are also appreciable changes in the APS with change in the sign of μ. The masses $m_{\tilde{\chi}^\pm}$ and $m_{\tilde{\tau}_L}$ increase significantly as one change $\mu < 0$ to $\mu > 0$. One need high value of $m_{1/2}$ to keep $m_{\tilde{\chi}^\pm}$ above experimental bound and high value of m_{16} for $m_{\tilde{\tau}_L}$ above experimental bound for $\mu < 0$. For example, $60(140) \lesssim m_{16} \lesssim 460(440)$ GeV for $\mu > 0(< 0)$, while the other parameters are as in Fig. 2.

We next examine the UFB3 constraint for the APS in Fig. 2. One of the important conclusions of this Letter is that the UFB-3 constraint rules out the entire APS for the universal model (throughout this Letter we shall use a $*$ $(+)$ for a UFB3 disallowed (allowed) points in the PS).

Next we will consider the effect of nonuniversality (compare Fig. 2 and Fig. 3). The SUSY parameters in Fig. 3 are as in Fig. 2 except that $m_{10} = 1.5\, m_{16}$. Such a modest non-universality may arguably appear

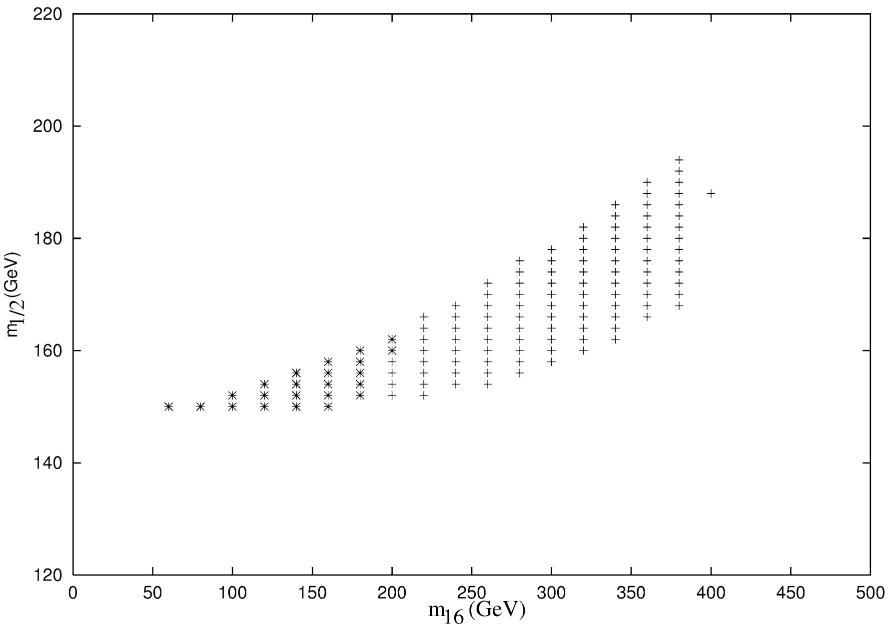

Fig. 3. The same as Fig. 2, with $m_{10} = 1.5 m_{16}$.

due to threshold corrections at M_G. For higher values of m_{10} μ^2 decreases rapidly and $m_{\tilde{\chi}^\pm}$ comes below experimental bound. A larger $m_{1/2}$ can avoid this problem but then the constraint $m_{\tilde{\nu}} < 80$ GeV requires a D-term that makes sfermion mass square negative at GUT scale. The overall APS, therefore, decreases. However, a region is still UFB3 allowed for $A_0 \gtrsim 0$, since $m^2_{H_u}$ is somewhat larger at M_G to begin with.

Next, we consider the possibility of Yukawa unification in this model [16]. It has already been shown in [10,17] that full t–b–τ Yukawa unification does not permit low slepton masses even in the presence of D-terms. We shall, therefore, restrict ourselves to partial b–τ unification with an accuracy $\leqslant 5\%$. We fix $\tan\beta$ to its lowest value required by unification. The APS in the universal model (Fig. 4) is qualitatively the same as in the fixed $\tan\beta$ case (compare Fig. 2 and Fig. 4) but its size somewhat smaller. It has been found that for higher values of m_D unification requires relatively low values of $\tan\beta \sim 20$. As indicated in Fig. 4 the APS is not consistent with the UFB3 constraint. Introduction of a modest non-universality at M_G as before, reduces the APS but leads to several UFB3 allowed points (Fig. 5). The following observations in the context of this model are worth noting. (i) We find a strong lower bounds $m_{\tilde{e}_R} \gtrsim 225$ GeV and $m_{d_R} \gtrsim 320$ GeV from the UFB3 constraint. (ii) We get a tight upper bound of $\tan\beta \lesssim 30$ independent of the choice of other parameters.

The phenomenological significance of a light sneutrino has already been discussed at length in the literature [18–24]. If the sneutrino mass happens to be in the range preferred by EWP data then it decays into the invisible channel $\tilde{\nu} \to \nu\chi_1^0$ with 100% BR and becomes an additional carrier of missing energy. If the lighter chargino mass happens to be near the current lower bound, a situation also preferred by EW precision data, then it would decay into the channel $\tilde{\chi}^\pm \to \ell\tilde{\nu}$ with almost 100% BR (the decay into sleptons are phase space suppressed), while in the conventional mSUGRA scenario it dominantly decay into jets. Finally, the second lightest neutralino $\tilde{\chi}_2^0$ which happens to be nearly degenerate with $\tilde{\chi}^\pm$ in models with gaugino mass unification, also decays dominantly into the invisible channel $\tilde{\chi}_2^0 \to \tilde{\chi}_1^0\tilde{\nu}$ and becomes another source of missing energy.

The additional carriers of missing energy which play roles similar to that of the lightest supersymmetric particle (LSP), may be termed virtual LSP (VLSP) in the context of collider experiments [18].

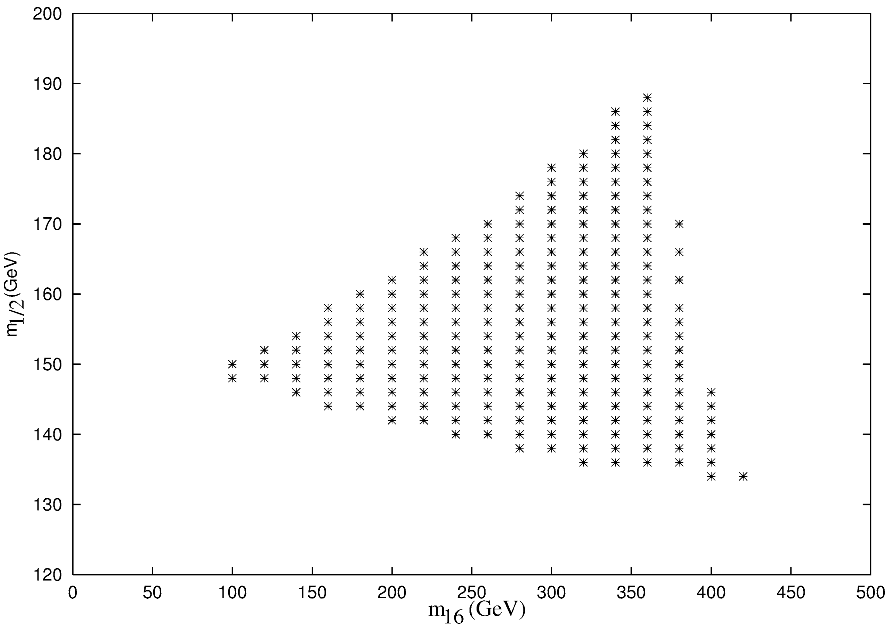

Fig. 4. The allowed parameter space in the universal scenario with b–τ unification. We set $A_0 = 0$ and m_D is fixed by the light slepton criterion. All points allowed by the Yukawa unification criterion are ruled out by UFB3.

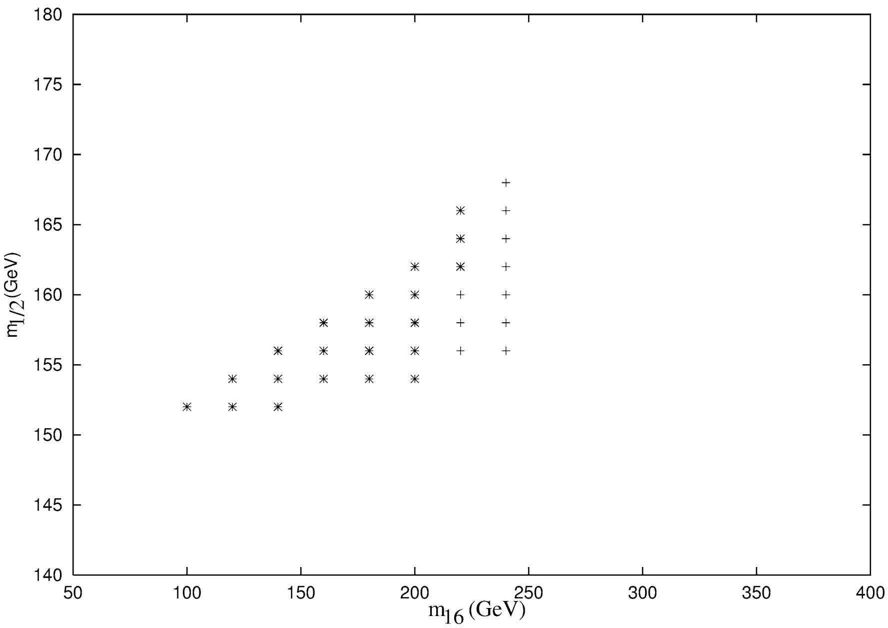

Fig. 5. The same as Fig. 4, with $m_{10} = 1.5\, m_{16}$.

In the VLSP scenario the collider signatures of squark–gluon production are quite different from the ones in conventional mSUGRA model due to the additional carriers of missing energy. Moreover, thanks to the enhanced leptonic decay of the chargino the lepton + jets + $\not{E}_T$ signal may increase at the cost of jets + $\not{E}_T$ signature [18,22]. The hadronically quiet tri-lepton signature [18] signalling the $\tilde{\chi}^{\pm}\tilde{\chi}_2^0$ production at the hadron colliders may disappear due to the invisible decay of $\tilde{\chi}_2^0$. On the other hand the hadronically quiet dilepton + $\not{E}_T$ signal from $\tilde{\chi}^{\pm}\tilde{\chi}^{\mp}$ may be boosted at the upgraded Tevatron as well at the e^+e^- colliders due to the enhanced leptonic decays of charginos [19,21]. Another dramatic signal of the VLSP model could be increase in the $e^+e^- \to \gamma+$ missing energy events [20]. In the conventional mSUGRA model the SUSY contributions comes only from the channel $e^+e^- \to \nu\tilde{\chi}_1^0\tilde{\chi}_1^0$ which has a modest cross section and is often swamped by the $e^+e^- \to \gamma\nu\bar{\nu}$ background. In the VLSP scenario, however, $e^+e^- \to \gamma\tilde{\nu}\tilde{\nu}^*, \gamma\tilde{\chi}_1^0\tilde{\chi}_2^0, \gamma\tilde{\chi}_2^0\tilde{\chi}_2^0$ contributes to the signal in addition to the above conventional mSUGRA process. Implementing some special cuts devised in [20] one can easily suppress the SM background. In particular, a suitable cut on the photon energy may kill a large number of $e^+e^- \to \gamma\nu\bar{\nu}$ events arising due to the radiative return to the Z peak at LEP energies above the Z pole without affecting the signal. A reanalysis of the LEP data using such cuts may reveal the VLSP scenario or severely restrict the sneutrino mass range preferred by EWP data.

If $m_{t_1} < m_{\tilde{\chi}^{\pm}}$, then the preferred decay mode of the lighter stop $(\tilde{t}_1)$ could be $\tilde{t} \to b\ell\tilde{\nu}$ rather the loop induced decay $\tilde{t} \to c\tilde{\chi}_1^0$ [22]. This would enhance the leptonic signal from the stop at the cost of jets + $\not{E}_T$ events.

While light sleptons may arise in many scenarios including the ones not based on supergravity (e.g., in the AMSB model), the simultaneous presence of relatively right down squarks and light sleptons would vindicate the $SO(10)$ D-term model. Enhancement of the jets + missing energy signal at the expenses of leptons + jets + $\not{E}_T$ signal from squark gluino production would be the hall-mark of this scenario [23, 25,26]. The effect becomes particularly striking if $m_{\tilde{g}} > m_{\tilde{d}_R}$, while all other squarks are much heavier than the gluinos [25,26]. This mass hierarchy is in fact obtained over the bulk of the parameter space probed in this Letter.

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